

Contents lists available at ScienceDirect

Physics Letters B

www.elsevier.com/locate/physletb

Scattering processes could distinguish Majorana from Dirac neutrinos



J. Barranco, D. Delepine*, V. Gonzalez-Macias, C. Lujan-Peschard, M. Napsuciale

Departamento de Física, División de Ciencias e Ingeniería, Campus León, Universidad de Guanajuato, León 37150, Mexico

ARTICLE INFO

Article history:

Received 24 August 2014

Received in revised form 2 November 2014

Accepted 7 November 2014

Available online 11 November 2014

Editor: A. Ringwald

ABSTRACT

It is well known that Majorana neutrinos have a pure axial neutral current interaction while Dirac neutrinos have the standard vector-axial interaction. In spite of this crucial difference, usually Dirac neutrino processes differ from Majorana processes by a term proportional to the neutrino mass, resulting in almost unmeasurable observations of this difference. In the present work we show that once the neutrino polarization evolution is considered, there are clear differences between Dirac and Majorana scattering on electrons. The change of polarization can be achieved in astrophysical environments with strong magnetic fields. Furthermore, we show that in the case of unpolarized neutrino scattering onto polarized electrons, this difference can be relevant even for large values of the neutrino energy.

© 2014 The Authors. Published by Elsevier B.V. This is an open access article under the CC BY license (<http://creativecommons.org/licenses/by/3.0/>). Funded by SCOAP³.

There are still many open questions in particle physics, and many of these involve the leptonic sector: what is the neutrino mass scale and the neutrino mass hierarchy? why is the neutrino mixing matrix so different from quarks? is there CP violation in the neutrino sector? Is the neutrino a Dirac or a Majorana particle? Pontecorvo called this last question the *central problem in neutrino physics*. If neutrino is a Majorana particle, then the neutrino is identical to its own antiparticle. If this is the case, the neutrinoless double beta decay is possible [1]. If such a process is experimentally observed, it will be an undoubted signal of the Majorana nature of the neutrino [2]. If, on the other hand, the neutrino is a Dirac particle, then the antineutrino is a different particle than the neutrino.

There is another crucial difference between Dirac and Majorana neutrinos [3–5]. If we consider the neutrino–electron scattering, either Dirac or Majorana, the effective Lagrangian at low energies can be written as:

$$\mathcal{L}_{ve} = \frac{G_F}{\sqrt{2}} [\bar{u}_{\nu\ell} \gamma^\mu (1 - \gamma^5) u_{\nu\ell}] [\bar{u}_e \gamma_\mu (g_V^\ell - g_A^\ell \gamma^5) u_e], \quad (1)$$

where the coupling constants are given by

$$g_V^\ell = -\frac{1}{2} + 2 \sin^2 \theta_W + \delta_{le}, \quad g_A^\ell = -\frac{1}{2} + \delta_{le}, \quad (2)$$

$$\ell = e, \mu, \tau.$$

The amplitude for the neutrino–electron scattering in the Dirac case is:

$$\mathcal{M}^D(\nu_\ell e \rightarrow \nu_\ell e) = -i \frac{G_F}{\sqrt{2}} [\bar{u}_e^f \gamma^\mu (g_V^\ell - g_A^\ell \gamma^5) u_e^i] [\bar{u}_{\nu\ell}^f \gamma_\mu (1 - \gamma^5) u_{\nu\ell}^i], \quad (3)$$

while in the Majorana case, since the neutrino is its own antiparticle, the amplitude will be:

$$\begin{aligned} \mathcal{M}^M(\nu_\ell e \rightarrow \nu_\ell e) &= -i \frac{G_F}{\sqrt{2}} [\bar{u}_e^f \gamma^\mu (g_V^\ell - g_A^\ell \gamma^5) u_e^i] \\ &\quad \times [\bar{u}_{\nu\ell}^f \gamma_\mu (1 - \gamma^5) u_{\nu\ell}^i - \bar{v}_{\nu\ell}^f \gamma_\mu (1 - \gamma^5) v_{\nu\ell}^i]. \end{aligned} \quad (4)$$

If the neutrino is a Majorana particle, then the following identity is valid:

$$\bar{v}_{\nu\ell}^f \gamma_\mu (1 - \gamma^5) v_{\nu\ell}^i = \bar{u}_{\nu\ell}^f \gamma_\mu (1 + \gamma^5) u_{\nu\ell}^i, \quad (5)$$

hence, the amplitude in the Majorana case will be

$$\begin{aligned} \mathcal{M}^M(\nu_\ell e \rightarrow \nu_\ell e) &= i \frac{2G_F}{\sqrt{2}} [\bar{u}_e^f \gamma^\mu (g_V^\ell - g_A^\ell \gamma^5) u_e^i] [\bar{u}_{\nu\ell}^f \gamma_\mu \gamma^5 u_{\nu\ell}^i]. \end{aligned} \quad (6)$$

It is clear that Eq. (3) is very different from Eq. (6). Nevertheless, the neutrino mass is extremely small ($m_\nu < 2$ eV [6]), thus these are almost completely chiral states, that is, almost fully polarized

* Corresponding author.

E-mail address: delepine@fisica.ugto.mx (D. Delepine).

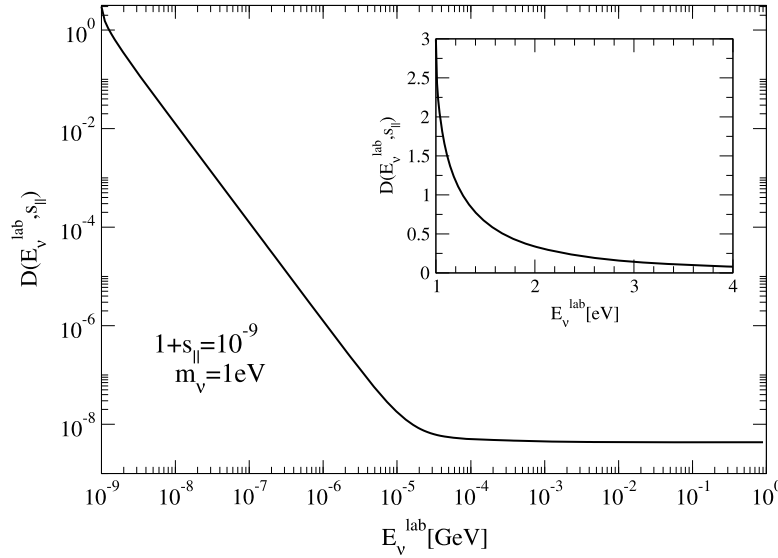


Fig. 1. Difference between the Majorana and Dirac neutrino–electron elastic scattering for a longitudinal polarization $1 + s_{\parallel} = 10^{-9}$ and a neutrino mass $m_{\nu} = 1$ eV. Inner figure shows a zoom for low neutrino energies.

particles due to the left-handed nature of the charged weak interaction. For this reason, an extra state preparation factor $(1 - \gamma_5)/2$ is usually added such that Eqs. (3) and (6) become identical [3]. The additional preparation factor is true if the neutrino mass is zero, but once the neutrino mass is incorporated, the neutrino is not completely polarized. Indeed, for instance, for the pure leptonic decay of a pseudoscalar meson $P^+ \rightarrow \ell^+ + \nu_{\ell}$, the neutrino longitudinal polarization is a function that depends on the neutrino mass [7]:

$$P_{\text{long}} = \frac{(E - W)|\vec{k}|}{WE - |\vec{k}|^2}, \quad (7)$$

with W and E the energies of the charged lepton ℓ and the neutrino, respectively. Here \vec{k} is given by $m_{\nu}^2 = m_p^2 + m_{\ell}^2 - 2m_p\sqrt{m_{\ell}^2 + |\vec{k}|^2}$, m_p the mass of the pseudoscalar meson and m_{ℓ} the lepton mass and m_{ν} the neutrino mass. For a neutrino mass of 1 eV, the polarization differs from a completely left-handed lepton in one part in a billion, thus, the prescription of adding a preparation factor seems reasonable. Nevertheless, since we are in the high precision test of the Standard Model era, it is reasonable to evaluate the differences between the Dirac and Majorana neutrino–electron scattering cross sections considering that neutrinos are highly polarized but not completely polarized. A straightforward calculation of the neutrino–electron scattering $\nu_{\ell}(p_{\nu}, s_{\nu}) + e(p_e) \rightarrow \nu_{\ell}(p'_{\nu}) + e(p'_e)$, with the incident neutrino polarization vector defined as $s_{\nu} = (0, s_{\perp}, 0, s_{\parallel})$ in the neutrino rest frame, gives in the Dirac case [3]:

$$\begin{aligned} \frac{d\sigma^D}{d\Omega} = & \frac{G_F^2}{8\pi^2 s} ((m_e^2(E_{\nu} - p^2 \cos \theta)(g_A^{\ell 2} - g_V^{\ell 2}) \\ & + (E_{\nu} E_e + p^2)(g_V^{\ell} + g_A^{\ell})^2 \\ & + (E_{\nu} E_e + p^2 \cos \theta)^2 (g_V^{\ell} - g_A^{\ell})^2 \\ & - p[s^{1/2}(E_{\nu} E_e + p^2)s_{\parallel}(g_V^{\ell} + g_A^{\ell})^2 + (E_{\nu} E_e + p^2 \cos \theta) \\ & \times ((E_e + E_{\nu} \cos \theta)s_{\parallel} + m_{\nu}s_{\perp} \sin \theta \cos \phi])(g_V^{\ell} - g_A^{\ell})^2 \\ & + m_e(E_{\nu}(1 - \cos \theta)s_{\parallel} - m_{\nu}s_{\perp} \sin \theta \cos \phi) \\ & \times (g_A^{\ell 2} - g_V^{\ell 2})), \end{aligned} \quad (8)$$

while in the Majorana case it is given by [3]:

$$\begin{aligned} \frac{d\sigma^M}{d\Omega} = & \frac{G_F^2}{4\pi^2 s} (((E_{\nu} E_e + p^2)^2 + (E_{\nu} E_e + p^2 \cos \theta)^2 \\ & + m_{\nu}(E_{\nu}^2 - p^2 \cos \theta))(g_V^{\ell 2} + g_A^{\ell 2}) \\ & + m_e^2(E_{\nu}^2 - p^2 \cos \theta + 2m_{\nu}^2)(g_A^{\ell 2} - g_V^{\ell 2}) \\ & - 2g_V^{\ell} g_A^{\ell} p(2E_{\nu} E_e + p^2(1 + \cos \theta)) \\ & \times (E_{\nu}s_{\parallel}(1 - \cos \theta) - m_{\nu}s_{\perp} \sin \theta \cos \phi)). \end{aligned} \quad (9)$$

In both cases, the variables $E_e, E_{\nu}, \theta, \phi$ and p refer to center of mass (CM) quantities and s the Mandelstam variable.

In order to quantify any difference between the Majorana and the Dirac cases, we define the function

$$D(E_{\nu}^{\text{lab}}, s_{\parallel}) = \frac{|\sigma(\nu_{\text{pol}}^D e) - \sigma(\nu_{\text{pol}}^M e)|}{\sigma(\nu_{\text{pol}}^D e)}, \quad (10)$$

where we have integrated Eqs. (8) and (9) over the CM angles and we have changed from the CM frame to the laboratory frame. This difference is shown in Fig. 1 and it summarizes a long discussion about the possibility of distinguishing Dirac from Majorana neutrinos in neutrino–electron scattering processes [3,8–10]: for terrestrial experiments where neutrinos are produced via charged currents, it is extremely difficult to observe significant differences between Dirac and Majorana neutrinos. Indeed, as can be seen in Fig. 1, for detectable neutrino energies, the difference is negligible. It becomes significant only for unreasonable (of the order of eV) energies of the neutrino.

Despite this fact, it is important to remember that neutrinos can forget its chiral origin. Indeed, any particle possessing a magnetic moment, as the neutrino does, interacts with external electromagnetic fields and consequently, its spin may rotate around the direction imposed by this external field. Furthermore, neutrinos can have a non-negligible magnetic moment μ_{ν} . Actually, current experimental constraints only gives a superior bound on the neutrino magnetic moment, $\mu_{\nu} < 3.2 \times 10^{-11} \mu_B$ [6] with μ_B the Bohr magneton. This limit is very big as compared with the expected neutrino magnetic moment that can arise from radiative corrections in the Standard Model. For example, for a Dirac neutrino $\mu_{\nu} \sim 3 \times 10^{-19} (\frac{m_{\nu}}{1 \text{ eV}}) \mu_B$ [11]. Hence, the spin of the neutrino could have a precession.

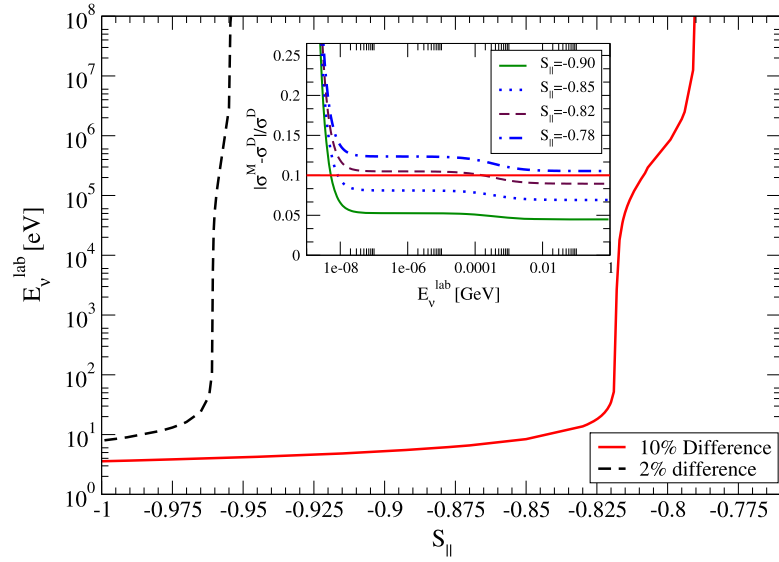


Fig. 2. Neutrino energy as a function of the neutrino longitudinal polarization needed to have: 10% difference between the Dirac and Majorana cross section (red solid line) and 2% (black dashed line).

In the past, this spin precession has been used as a mechanism to probe the Majorana nature of the neutrino [12,13], but those works have focused on a complete transition from a Majorana neutrino to an antineutrino due to the spin precession. The non-observation of solar antineutrinos in KamLAND has derived an upper limit on the neutrino magnetic moment only $\mu_\nu < 5 \times 10^{-12} \mu_B$ [14].

We will show, that contrary to those previous attempts, it is not necessary to have a completely flipped neutrino, that is an antineutrino, to have observable differences between the Dirac and Majorana neutrinos. This can be observed by re-evaluating Eq. (10) for different values of the neutrino longitudinal polarization $s_{||}$. Inner plot of Fig. 2 shows the function $D(E_\nu^{\text{lab}}, s_{||})$ for different values of $s_{||}$. As it can be seen, the value of the neutrino energy where a significant difference between Dirac and Majorana scattering cross section appears is a function of $s_{||}$. In Fig. 2 we show the value of the neutrino energy for two cases: when the difference is 10% (solid red line) and 2% (dashed black line). As it can be seen, there is an asymptotic value of $s_{||}$ where this difference is reachable for current neutrino detectors. In the first case, it is needed that the neutrino forgets its chiral origin almost 25%, while a more precise experiment able to detect astrophysical neutrinos with a 2% accuracy will need only a 5% deviation in the neutrino's original helicity.

What is the magnetic field needed in order to have such changes in the neutrino's helicity? In order to estimate this, we recall previous studies where the depolarization rate of neutrinos was calculated [15–17]. In the case of a random distribution of electromagnetic fields, the average neutrino's helicity $\langle h \rangle$ changes as dictated by the equation $\langle h(t) \rangle = \exp(-\Gamma_{\text{depol}} t) \langle h(0) \rangle$, where

$$\Gamma_{\text{depol}} = 0.0132 \mu_\nu^2 T^3, \quad (11)$$

μ_ν the magnetic moment of the neutrino [17]. Another source for neutrino spin depolarization is produced by the interaction with a large scale magnetic field. For this case, the depolarization rate is given by

$$\Gamma_B = \frac{4\mu_\nu^2 B^2 D}{\omega_{\text{refr}}^2 + 4\mu_\nu^2 B^2}, \quad (12)$$

where B is the external magnetic field. For electron neutrinos, $\omega_{\text{refr}} \simeq 1.1 \times 10^{-20} (T/\text{MeV})^4 \langle p_0 \rangle$, p_0 the neutrino energy, and $D \simeq 2.04 G_F^2 T^5$ [17].

Let us focus in the case of a random distribution of magnetic fields, Eq. (11). In Fig. 3 we have plotted the depolarization distance $\lambda_{\text{depol}} = 1/\Gamma_{\text{depol}}$ as a function of the neutrino magnetic moment. Solid black line is for a complete depolarization while red dashed line for a 5% change in the average neutrino helicity. This is for a temperature of $T = 20$ MeV. Dotted lines represent typical astrophysical objects. The solar radius has been plotted just as a reference distance. The temperature in the core of the Sun is just a few keV, and then Eq. (11) is not valid for those low temperatures.

Moreover, Fig. 3 implies that other neutrinos produced in different astrophysical objects could in principle be depolarized if the neutrino magnetic moment is relatively large compared to the SM prediction.

Since neutrinos can have a broad distribution of spin polarization, finally, let us consider the extreme case where neutrinos are unpolarized and compute the neutrino–electron elastic scattering. As previously noted [10], in this case the corresponding matrix elements for Dirac and Majorana neutrinos are completely different. Furthermore, as Eqs. (3) and (6) show, the difference increments as long as we maximize the axial contribution. In order to do that, we will consider that neutrinos are unpolarized and that the electron in the target is polarized. The possible use of polarized electrons in ν – e elastic scattering has been previously proposed in order to look for other non-standard neutrino interactions, as the neutrino magnetic moment itself [18], or possible CP violation signals [19] among other different motivations [20]. An efficient Polarized Electron Target (PET) can be obtained applying strong, although realistic, magnetic fields to the detector [18,19]. Putting aside for the moment the technical difficulties in order to have such PET detector, here we will show that the electron polarization could enhance the difference between Dirac and Majorana neutrinos as long as neutrinos had lost their initial polarization due to an interaction of the neutrino with external magnetic fields in some wild astrophysical environment.

To take into account the electron polarization, we used the Michel–Wightman [21] formalism. For the evaluations we used the laboratory frame, in which the electron is at rest and the electron polarization vector angle ξ is given respect to the direction of

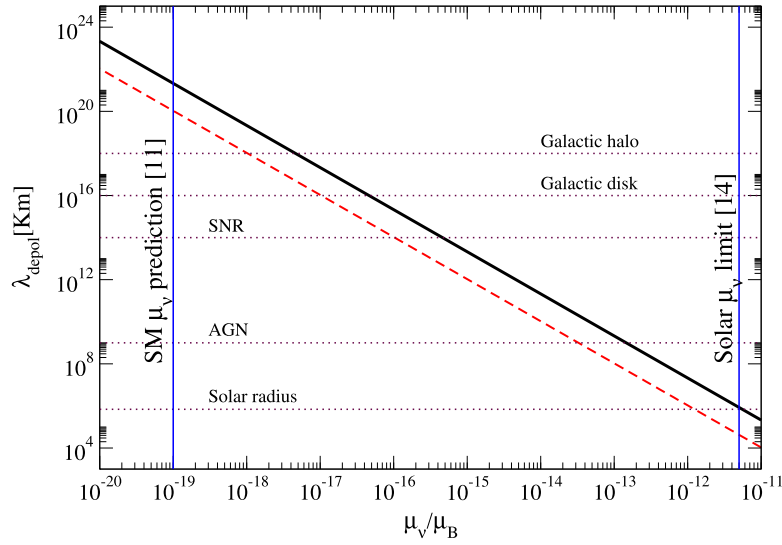


Fig. 3. Depolarization distance for a random distribution of magnetic fields as a function of the neutrino magnetic moment. Solid black line is for a complete depolarization while red dashed line for a 5% change in the average neutrino helicity. Dotted lines represents typical astrophysical objects: The solar radius, the size of an Active Galactic Nuclei (AGN), a Super Nova Remnant (SNR), the size of the galaxy disk and the galaxy halo. We have assumed $T = 20$ MeV. Solar radius is given just as a reference distance, in case a hypothetical hot object with such radius could emit large amount of neutrinos.

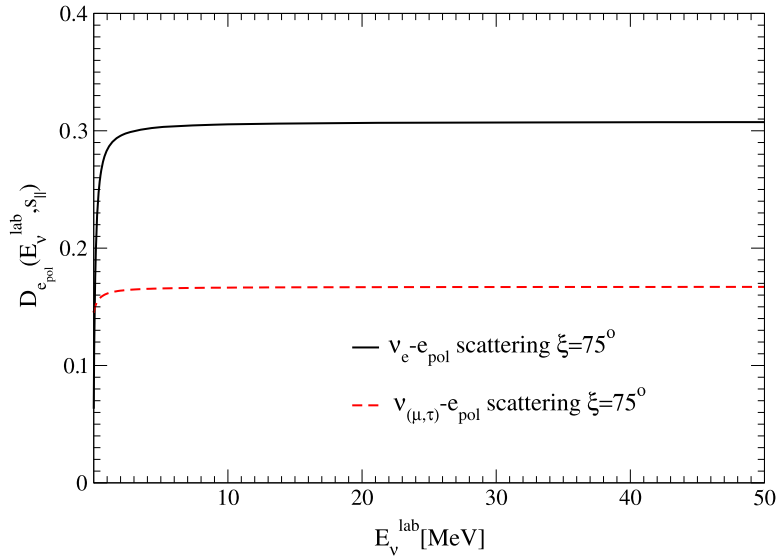


Fig. 4. Difference for the neutrino–polarized electron scattering for a fixed value of the target electron polarization as defined by Eq. (15).

the incoming neutrino (that we choose to be z): $s_e = (0, s_\perp, 0, s_\parallel)$, $s_\parallel = \cos \xi$.

In the following we give the expressions for the differential cross section for the elastic scattering of $\nu_\ell - e$ considering the target electrons can be polarized in any direction. We will keep the dependence on the angle $s_\parallel = \cos \xi$ with respect to the incoming neutrinos direction explicitly.

We have neglected the neutrino mass, since the changes in the cross sections due to the inclusion of the neutrino mass are very small.

We arrive to the Dirac neutrino–polarized electron elastic scattering:

$$\frac{d\sigma(\nu_\ell^D e_{\text{pol}})}{dT_e} = \frac{G_F^2 m_e}{2\pi} \left[(g_A^\ell - g_V^\ell)^2 (1 - s_\parallel) \left(1 - \frac{T_e}{E_v^{\text{lab}}}\right)^2 + (1 + s_\parallel) \left((g_A^\ell + g_V^\ell)^2 + (g_A^{\ell 2} - g_V^{\ell 2}) \frac{m_e T_e}{E_v^{\text{lab} 2}} \right) \right]$$

$$+ (g_A^\ell - g_V^\ell)^2 s_\parallel \left(1 - \frac{T_e}{E_v^{\text{lab}}}\right) \frac{m_e T_e}{E_v^{\text{lab} 2}} \right]. \quad (13)$$

Here, we have chosen the laboratory frame, with T_e the electron recoil energy as it was done in [18], and our results match. Recall g_A^ℓ, g_V^ℓ are defined in Eq. (2). On the other hand, in the Majorana case, the neutrino–polarized electron cross sections is given by:

$$\begin{aligned} \frac{d\sigma(\nu_\ell^M e_{\text{pol}})}{dT_e} &= \frac{G_F^2 m_e}{2\pi} \left[2(g_A^{\ell 2} + g_V^{\ell 2}) \left(1 - \frac{T_e}{E_v^{\text{lab}}}\right)^2 + 2(g_A^{\ell 2} + g_V^{\ell 2}) + 2(g_A^{\ell 2} - g_V^{\ell 2}) \frac{m_e T_e}{E_v^{\text{lab} 2}} + 4g_A^\ell g_V^\ell s_\parallel \left(1 + \left(1 - \frac{T_e}{E_v^{\text{lab}}}\right)^2 - \left(1 - \frac{T_e}{E_v^{\text{lab}}}\right) \frac{m_e T_e}{E_v^{\text{lab} 2}} \right) \right]. \quad (14) \end{aligned}$$

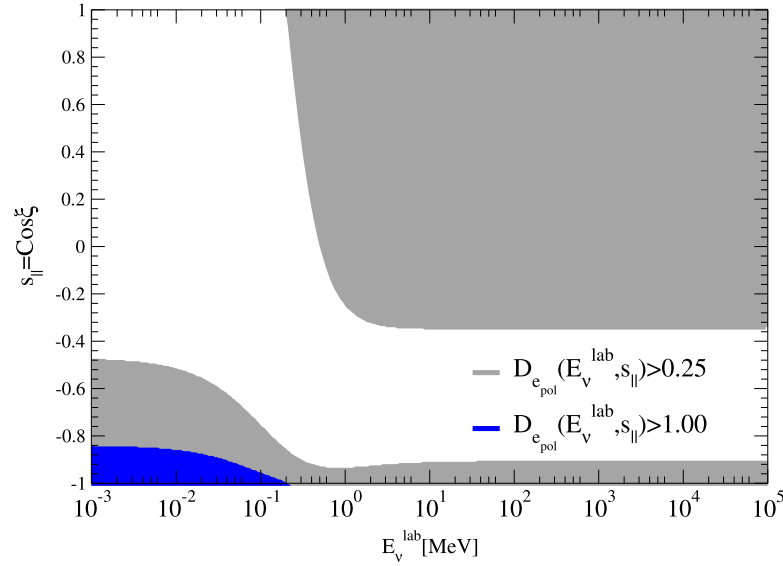


Fig. 5. Electron polarization needed to have significant differences between Dirac and Majorana cross sections for unpolarized neutrinos as a function of the incoming neutrino energy.

After an integration over the electron recoil energy, $0 < T_e < \frac{2E_{\nu}^{\text{lab}}}{m_e + 2E_{\nu}^{\text{lab}}}$ we can compute the difference:

$$D_{e\text{pol}}(E_{\nu}^{\text{lab}}, s_{\parallel}) = \frac{|\sigma(\nu^D e_{\text{pol}}) - \sigma(\nu^M e_{\text{pol}})|}{\sigma(\nu^D e_{\text{pol}})}, \quad (15)$$

for different values of the electron polarization. An explicit example is shown in Fig. 4.

It is remarkable that an appreciable difference is obtained for detectable energies of the incoming neutrino.

Actually, in the case of unpolarized neutrinos that scatter on polarized electrons there could be differences as big as twice the cross section, i.e. $D_{e\text{pol}}(E_{\nu}^{\text{lab}}, s_{\parallel}) > 1$, for certain values of the neutrino energy and the degree of polarization of the target electrons. This is illustrated in Fig. 5, where we have plotted isocurves of $D_{e\text{pol}}(E_{\nu}^{\text{lab}}, s_{\parallel})$. Although this extreme case is reachable only for extremely low energetic neutrinos. Nevertheless, the case $D_{e\text{pol}}(E_{\nu}^{\text{lab}}, s_{\parallel}) > 0.25$, i.e. differences of at least 25% are expected for a wide range in the electron polarization and neutrino energy.

In summary, we have shown that once the neutrino polarization evolution is considered, there are clear differences between Dirac and Majorana scattering on electrons. This change in the evolution of the helicity is possible due to the existence of a neutrino magnetic moment. The change of polarization can be achieved in astrophysical environments with magnetic fields if the neutrino magnetic moment is bigger than the expected SM prediction but smaller than the current experimental limit. This helicity change could be strong enough to completely unpolarize the neutrino. Furthermore, we show that in the case of unpolarized neutrino scattering onto polarized electrons, this difference can be relevant even for larger values of the neutrino energy.

Acknowledgements

This work has been supported by CONACyT SNI-Mexico. The authors are also grateful to Conacyt (CB-156618), DAIP project (Guanajuato University) and PIFI (SEP) for financial support.

References

- [1] W.H. Furry, Phys. Rev. 56 (1939) 1184.
- [2] J. Schechter, J.W.F. Valle, Phys. Rev. D 25 (1982) 2951.
- [3] B. Kayser, R.E. Shrock, Phys. Lett. B 112 (1982) 137.
- [4] R.N. Mohapatra, P.B. Pal, World Sci. Lect. Notes Phys. 60 (1998) 1, World Sci. Lect. Notes Phys. 72 (2004) 1.
- [5] J.-E. Campagne, LAL-95-23.
- [6] J. Beringer, et al., Particle Data Group, Phys. Rev. D 86 (2012) 010001.
- [7] G. Barenboim, J. Bernabeu, O. Vives, Phys. Rev. Lett. 77 (1996) 3299, arXiv:hep-ph/9606218.
- [8] G.V. Dass, Phys. Rev. D 32 (1985) 1239.
- [9] T. Garavaglia, Phys. Rev. D 29 (1984) 387.
- [10] S. Hannestad, arXiv:hep-ph/9701216.
- [11] K. Fujikawa, R. Shrock, Phys. Rev. Lett. 45 (1980) 963.
- [12] V.B. Semikoz, Nucl. Phys. B 498 (1997) 39, arXiv:hep-ph/9611383.
- [13] S. Pastor, V.B. Semikoz, J.W.F. Valle, Phys. Lett. B 423 (1998) 118, arXiv:hep-ph/9711316.
- [14] O.G. Miranda, T.I. Rashba, A.I. Rez, J.W.F. Valle, Phys. Rev. Lett. 93 (2004) 051304, arXiv:hep-ph/0311014.
- [15] A. Loeb, L. Stodolsky, Phys. Rev. D 40 (1989) 3520.
- [16] V. Semikoz, Phys. Rev. D 48 (1993) 5264; V. Semikoz, Phys. Rev. D 49 (1994) 6246 (Erratum).
- [17] P. Elmfors, K. Enqvist, G. Raffelt, G. Sigl, Nucl. Phys. B 503 (1997) 3, arXiv:hep-ph/9703214.
- [18] T.I. Rashba, V.B. Semikoz, Phys. Lett. B 479 (2000) 218, arXiv:hep-ph/0003099.
- [19] S. Ciechanowicz, W. Sobkow, M. Misiaszek, Phys. Rev. D 71 (2005) 093006, arXiv:hep-ph/0309286.
- [20] P. Minkowski, M. Passera, Phys. Lett. B 541 (2002) 151, arXiv:hep-ph/0201239.
- [21] L. Michel, A.S. Wightman, Bulletin of the American Physical Society 30 (1) (1955) 1190; M. Kaku, Quantum Field Theory, Oxford Univ. Pr., New York, USA, 1993, p. 84, 785 pp.